

# Learning from the New Higgs-like Scalar before It Vanishes

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Motivated by a di-photon anomaly observed by ATLAS and CMS we develop an SFITTER analysis for a combined electroweak-Higgs sector, and a scalar portal at the LHC. The theoretical description is based on the linear effective Lagrangian for the Higgs and gauge fields, combined with an additional singlet scalar. The key feature is the extraction of reliable information on the portal structure of the combined scalar potential. For the specific di-photon anomaly we find that the new state might well form such a Higgs portal. To obtain more conclusive results we define and test the connection of the Wilson coefficients in the Higgs and heavy scalar sectors, as suggested by a portal setup.

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## I. INTRODUCTION

The discovery of a light Higgs boson [1, 2] has opened a major new avenue in experimental and theoretical particle physics: comprehensive tests of a possible non-minimal fundamental scalar sector, for which there exists a plethora of motivations. While there has been a lot of progress in developing combined Higgs and gauge analysis strategies for the LHC Run II [3–6], there exists no general and proven analysis framework even for a Higgs portal model [7].

The announcement of an excess seen in the di-photon spectrum by both ATLAS and CMS [8–10], if confirmed by future data, suggests such an extended scalar sector. The anomaly has led to an excessive number of publications, so we feel that adding one more, and hopefully useful publication can be justified somehow\*. Early studies of the anomaly in an effective theory framework can be found in Ref. [11]. Intriguingly, an additional scalar is not sufficient to explain the signal in complete models. The new scalar’s sizable couplings to photons and gluons need to be induced by relatively light new particles [12]. For example in supersymmetric models, vector-like matter added to the MSSM or non-trivial signatures in the NMSSM are necessary for a successful explanation of the excess [13]. In models in which these new states are connected to the SUSY breaking sector the new scalar can be identified with the sgoldstino, implying a very low SUSY breaking scale [14]. Other extended spacetime symmetries give rise to dilaton [15] and radion interpretations [16], which imply similarly unintended consequences, such as low ultraviolet (UV) scales or a very large curvature of the extra dimension. Extra dimensional scalars not directly related to the compactification circumvent this problem [17] and can explain the localization of extra dimensional fermions, which makes the new scalar a localizer field [18]. Related models, which consider the electroweak scale (or the TeV scale) arising from composite dynamics are less constrained than the MSSM, due to the large number of potential scalar resonances and fermionic quark partners [19]. The possibility of the new resonance to be a spin 2 particle, associated with a higher dimensional theory of gravity is strongly constrained by dilepton searches and just like the radion implies sizable curvature terms [20]. The large width of  $\Gamma_S = 45$  GeV, as reported by ATLAS, can be addressed in some models [21, 22]; while such a large width only slightly increases the statistical significance, if it is true, background interference effects are important [23]. In this case, it is well motivated to assume that the new scalar provides a portal to a dark sector, inducing a sizable width through invisible decays [24]. Alternatively, it might be the sign of cascade decays or other explanations not based on a single scalar resonance, which lead to cusps and endpoint structures that can fake a large width [25]. The very minimal, yet not perturbatively realizable assumption of photon fusion induced production can not explain a large width [26]. The new resonance could also be related to the various, persistent flavor anomalies [27], to the mechanism behind the electroweak phase transition [28], the strong CP problem [29], or an underlying string theory [30]. Finally, a variety of models, motivated by different extensions of the Standard Model (SM) not fitting in the above categories, and further measurements testing the properties of the new resonance have been proposed [31].

In spite of all these considerations, the most obvious question is whether such an additional, likely scalar resonance can be part of an extended Higgs sector [32, 33]; in other words, if the new scalar can form a Higgs portal, possibly to a new sector. To answer this question we will remain agnostic about the underlying physics, but assume that a resonantly produced narrow scalar singlet is responsible for the excess. We couple the new scalar to the SM through an effective Lagrangian. This assumption exactly corresponds to recent developments on how to describe deviations from the Standard Model Higgs sector at the LHC [3–6, 34–37]. The combined Higgs portal Lagrangian

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\* Beyond the di-photon anomaly we present the first full SFITTER analysis of a Higgs portal allowing for higher-dimensional operators.

is organized by the field content, the symmetry structure, and the mass dimension. This way we can contrast the apparent absence of dimension–six effects in the range  $\Lambda \approx 300 \dots 500$  GeV for the SM-like Higgs and gauge sector [5] with the need for higher-dimensional operators coupling to the new scalar with  $\Lambda \lesssim 1$  TeV.

### A. Theoretical framework

The most general linear effective Lagrangian up to dimension six and built from Standard Model particles and a new scalar singlet reads

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_{\text{dim-6}}^H + \mathcal{L}_{\text{dim}\leq 5}^S + \mathcal{L}_{\text{dim-6}}^S. \quad (1)$$

Here,  $\mathcal{L}_{\text{SM}}$  stands for the renormalizable SM Lagrangian, while  $\mathcal{L}_{\text{dim-6}}^H$  contains the dimension–six operators made out of SM fields. Adopting the basis of our set of Higgs legacy papers [3–5] it reads

$$\begin{aligned} \mathcal{L}_{\text{dim-6}}^H = & \frac{f_{BB}}{\Lambda^2} \phi^\dagger \hat{B}_{\mu\nu} \hat{B}^{\mu\nu} \phi + \frac{f_{WW}}{\Lambda^2} \phi^\dagger \hat{W}_{\mu\nu} \hat{W}^{\mu\nu} \phi - \frac{\alpha_s}{8\pi} \frac{f_{GG}}{\Lambda^2} \phi^\dagger \phi G_{\mu\nu}^a G^{a\mu\nu} + \frac{f_{WWW}}{\Lambda^2} \text{tr} \left( \hat{W}_{\mu\nu} \hat{W}^{\nu\rho} \hat{W}_\rho{}^\mu \right) \\ & + \frac{f_B}{\Lambda^2} (D_\mu \phi)^\dagger \hat{B}^{\mu\nu} (D_\nu \phi) + \frac{f_W}{\Lambda^2} (D_\mu \phi)^\dagger \hat{W}^{\mu\nu} (D_\nu \phi) + \frac{f_{\phi,2}}{\Lambda^2} \frac{1}{2} \partial^\mu (\phi^\dagger \phi) \partial_\mu (\phi^\dagger \phi) \\ & + \left( \frac{f_\tau m_\tau}{v \Lambda^2} (\phi^\dagger \phi) (\bar{L}_3 \phi e_{R,3}) + \frac{f_b m_b}{v \Lambda^2} (\phi^\dagger \phi) (\bar{Q}_3 \phi d_{R,3}) + \frac{f_t m_t}{v \Lambda^2} (\phi^\dagger \phi) (\bar{Q}_3 \phi u_{R,3}) + \text{h.c.} \right). \quad (2) \end{aligned}$$

The Higgs covariant derivative is  $D_\mu \phi = (\partial_\mu + ig' B_\mu/2 + ig \sigma_a W_\mu^a/2) \phi$ , and the field strengths are  $\hat{B}_{\mu\nu} = ig' B_{\mu\nu}/2$  and  $\hat{W}_{\mu\nu} = ig \sigma^a W_{\mu\nu}^a/2$  in terms of the Pauli matrices  $\sigma^a$ . The  $SU(2)_L$  and  $U(1)_Y$  gauge couplings are  $g$  and  $g'$ , respectively. While the minimum independent set consists of 59 baryon number conserving operators, barring flavor structure and Hermitian conjugation [37], we follow the definition of the relevant operator basis describing Higgs coupling and triple gauge boson vertex (TGV) modifications at the LHC in Ref. [3]. In our construction we assume a narrow, CP-even Higgs, focusing on the minimal, Yukawa-like, couplings to the heavy fermions. We use the equations of motion to rotate to a basis where there are no blind directions linked to electroweak precision data. That way, we neglect all operators contributing to electroweak precision observables at tree level in our LHC analysis. For the Standard Model fit [3–5] we omit the operator  $(\phi^\dagger \phi)^3$ , which only contributes to the rather poorly measured triple Higgs coupling. In the appendix we argue why even in the presence of an additional, mixing scalar, this operator will not add any extra relevant features to the fit.

Moving to the new scalar Lagrangian terms, we assume in the following that the additional singlet does not develop a VEV, or that the Lagrangian can be re-defined such that the VEV vanishes [33]. The effective Lagrangian of such an additional singlet scalar can be divided into two pieces. Following Refs. [38–41] we first write down a set of non-redundant, independent operators up to dimension five,

$$\begin{aligned} \mathcal{L}_{\text{dim}\leq 5}^S = & \frac{1}{2} \partial_\mu S \partial^\mu S - a_1 S - \frac{M_S^2}{2} S^2 - a_3 S^3 - a_4 S^4 - \frac{f_5^S}{\Lambda} S^5 \\ & - \mu_S S \phi^\dagger \phi - \frac{\lambda_{SH}}{2} S^2 \phi^\dagger \phi - \frac{f_1^S}{\Lambda} S (\phi^\dagger \phi)^2 - \frac{f_3^S}{\Lambda} S^3 \phi^\dagger \phi \\ & + \frac{\alpha_s}{4\pi} \frac{f_{GG}^S}{\Lambda} S G_{\mu\nu}^a G^{a\mu\nu} + \frac{\alpha}{4\pi c_w^2} \frac{f_{BB}^S}{\Lambda} S B_{\mu\nu} B^{\mu\nu} + \frac{\alpha}{4\pi s_w^2} \frac{f_{WW}^S}{\Lambda} S W_{\mu\nu}^a W^{a\mu\nu} \\ & \left( -\frac{f_d^S}{\Lambda} S \bar{Q}_L \phi d_R - \frac{f_u^S}{\Lambda} S \bar{Q}_L \tilde{\phi} u_R - \frac{f_\ell^S}{\Lambda} S \bar{L}_L \phi \ell_R + \text{h.c.} \right). \quad (3) \end{aligned}$$

To be fully consistent with the Standard Model Lagrangian we could then add all dimension–six operators including at least one power of the new singlet scalar. The corresponding set of additional operators can be written as [38]

$$\begin{aligned} \mathcal{L}_{\text{dim-6}}^S = & \frac{f_\phi^{SS}}{\Lambda^2} \phi^\dagger \phi \partial_\mu S \partial^\mu S - \frac{f_6^S}{\Lambda^2} S^6 - \frac{f_4^S}{\Lambda^2} S^4 \phi^\dagger \phi - \frac{f_2^S}{\Lambda^2} S^2 (\phi^\dagger \phi)^2 \\ & + \frac{f_{GG}^{SS}}{\Lambda^2} S^2 G_{\mu\nu}^a G^{a\mu\nu} + \frac{f_{BB}^{SS}}{\Lambda^2} S^2 B_{\mu\nu} B^{\mu\nu} + \frac{f_{WW}^{SS}}{\Lambda^2} S^2 W_{\mu\nu}^a W^{a\mu\nu} \\ & \left( -\frac{f_d^{SS}}{\Lambda^2} S^2 \bar{Q}_L \phi d_R - \frac{f_u^{SS}}{\Lambda^2} S^2 \bar{Q}_L \tilde{\phi} u_R - \frac{f_\ell^{SS}}{\Lambda^2} S^2 \bar{L}_L \phi \ell_R + \text{h.c.} \right) . \end{aligned} \quad (4)$$

Nevertheless, given the singlet nature of the new scalar and neglecting lepton number violation, all dimension–six operators including the singlet are quadratic in the field  $S$ . Consequently, their phenomenological effects will be contributions to the mass terms ( $f_2^S/\Lambda^2$ ), re-definitions of the  $S$  field to recover canonical kinetic terms ( $f_\phi^{SS}/\Lambda^2$ ), and the contributions to several vertices including two or more heavy scalars. After scalar-Higgs mixing, the two operators  $f_\phi^{SS}/\Lambda^2$  and  $f_2^S/\Lambda^2$  will contribute to the  $SHH$  interaction as well. However, all these phenomenology features are already taken into account in our analysis by the free parameters in the dimension–five Lagrangian. Therefore, we neglect for the time being the explicit features induced by Eq.(4). We give more details on the effective Lagrangian and the Higgs portal mixing in the Appendix.

## B. Analysis framework

The set of analyses presented here are derived using the SFITTER framework. SFITTER allows us to study multi-dimensional parameter spaces in the Higgs sector [4, 42], the gauge sector [5] and in new physics models like supersymmetry [43]. The fit procedure uses Markov chains to create an exclusive, multidimensional log-likelihood map, based on the available measurements and including all the relevant uncertainties and correlations. The construction of a profile likelihood with flat theory uncertainties leads to the RFIT scheme [44]. The statistic uncertainties on the measurements, both for event rates and kinematic distributions, follow Poisson statistics, as do the background uncertainties. All systematic uncertainties are described by Gaussian distributions and can be correlated between the relevant channels. We show log-likelihood projections on two-dimensional planes after profiling over all other parameters. Here, red-yellow regions will illustrate points within  $\Delta(-2\log \mathcal{L}) = 2.3$  of the best fit point log-likelihood ( $1\sigma$  in the Gaussian approximation), green regions indicate  $\Delta(-2\log \mathcal{L}) = 6.18$  ( $2\sigma$  in the Gaussian limit), and black dots imply the  $\Delta(-2\log \mathcal{L}) = 5.99$  exclusion limits (95% CL in the Gaussian case).

The implementation of experimental results in the SFITTER framework is described in Ref. [4] for the Higgs measurements and in Ref. [5] for anomalous triple gauge boson coupling measurements. For the triple gauge boson vertex (TGV) analyses<sup>†</sup> the correlation of the theory uncertainties between the different bins of a given kinematic distribution is taken into account by flat profiled nuisance parameters [5], while for the different Higgs channels the theory uncertainties are considered uncorrelated without a sizable impact on the shown results [4]. For the Higgs portal analysis we take into account the constraints on a possible new resonance based on the data listed in Tab. I. For the new resonance we only implement inclusive measurements assuming a narrow width.

<sup>†</sup> Note that pair production of weak bosons at the LHC is a crucial ingredient to a Higgs fit based on an effective Lagrangian assuming a linear realization of electroweak symmetry breaking. Without taking these measurements into account the qualitative and quantitative outcome of the fit will be wrong [5].

## II. HIGGS PORTAL ANALYSIS

In the following we will use the SFITTER effective Lagrangian framework to analyze a new gluon-fusion produced resonance in combination with the electroweak gauge and Higgs sectors at the weak scale. In other words, we ask the question whether such a new particle could be part of an extended Higgs sector and what the allowed parameter space is. In Sec. II A we only include the dimension-five operators given in Eq.(3), restricting the analysis to the data in Tab. I. In Sec. II B we combine this analysis with the Higgs-electroweak measurements and the SFITTER results induced by the dimension-six Lagrangian in Eq.(2). Finally we link the size of different operators to a common origin in Sec. II C.

### A. Heavy scalar fit

As a first step we analyze only the measurements for the heavy scalar, as listed in Tab. I. In Fig. 1 we use this data to determine the five parameters

$$\{ f_{WW}^S, f_{BB}^S, f_{GG}^S, \sin \alpha, c_{SHH} \} . \quad (5)$$

In our parametrization  $c_{SHH}$  accounts for the independent contributions to the  $SHH$  vertex from the dimension-five Lagrangian terms beyond the terms generating the mixing, as discussed in the Appendix.

The best fit point for this analysis has  $-2 \log L = 8.9$ , while the SM point leads to  $-2 \log L = 28.2$ , within a  $3.1\sigma$  range for a 5-parameter study (in the Gaussian limit). In the upper left panel of Fig. 1 we can see that within the displayed range of parameters both  $f_{WW}^S$  and  $f_{BB}^S$  are strongly correlated, and they present a flat direction. The correlation reflects the fact that they are the only Wilson coefficients contributing to the di-photon decay of the new scalar at tree level. Because  $f_{WW}^S$  is constrained through the decay  $S \rightarrow WW$ , the di-photon excess cannot be accommodated through this coupling only. Due to that we find that  $|f_{BB}^S/\Lambda| > 2 \text{ TeV}^{-1}$  in the upper-center panel. This is caused by the fact that  $f_{BB}^S$  does not contribute to the  $SWW$  vertex, and in addition its contribution to the  $SZZ$  vertex is suppressed by the weak mixing angle. This allows us to explain

Channel	Dataset	Reference
$S \rightarrow \gamma\gamma$	ATLAS 8 TeV	[9]
$S \rightarrow \gamma\gamma$	ATLAS 13 TeV	[9]
$S \rightarrow \gamma\gamma$	CMS 8 TeV	[45]
$S \rightarrow \gamma\gamma$	CMS 13 TeV	[10]
$S \rightarrow WW$	ATLAS 8 TeV	[46]
$S \rightarrow WW$	ATLAS 13 TeV	[47]
$S \rightarrow ZZ$	ATLAS 8 TeV	[48]
$S \rightarrow ZZ$	ATLAS 13 TeV	[49]
$S \rightarrow ZZ$	ATLAS 13 TeV	[50]
$S \rightarrow Z\gamma$	ATLAS 13 TeV	[51]
$S \rightarrow Z\gamma$	CMS 13 TeV	[52]
$S \rightarrow Z\gamma$	ATLAS 8 TeV	[53]
$S \rightarrow t\bar{t}$	ATLAS 8 TeV	[54]
$S \rightarrow jj$	CMS 8 TeV	[55]
$S \rightarrow hh$	ATLAS 13 TeV	[56]
$S \rightarrow hh$	CMS 8 TeV	[57]
$S \rightarrow \tau\bar{\tau}$	CMS 8 TeV	[58]

Table I. Experimental data on the heavy resonance included in our fit.

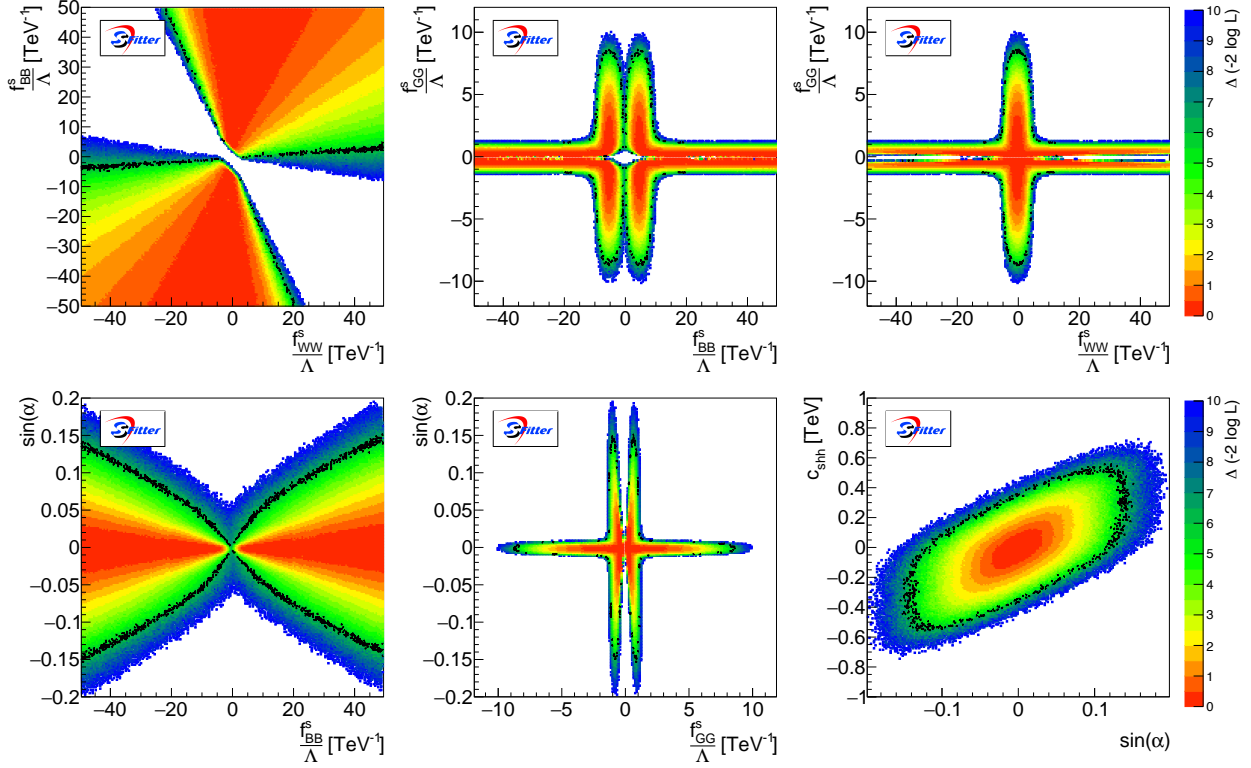


Figure 1. Two-dimensional profile log-likelihoods for the analysis of the heavy scalar sector alone spanning  $f_{WW}^S$ ,  $f_{BB}^S$ ,  $f_{GG}^S$ ,  $\sin \alpha$ , and  $c_{SHH}$ . The black points indicate  $\Delta(-2 \log L) = 5.99$ .

the observed excess without getting into conflict with the exclusion bounds, what makes  $f_{WW}^S = 0$  compatible with the best fit point, as shown in the upper-right panel.

Moving on to the mixing angle, we find  $\sin \alpha < 0.15$  at 95% CL and for the displayed ranges of  $f_{BB}^S$  and  $f_{WW}^S$  in the lower-left panel. This bound comes from the absence of a heavy scalar signal in  $WW$  and  $ZZ$ , but also in di-jet,  $t\bar{t}$ ,  $\tau\bar{\tau}$ , and  $hh$  decay channels. It is linked to maximum assumed values for  $f_{BB}^S$  and  $f_{WW}^S$ , because a larger mixing angle can be partially compensated by larger Wilson coefficients  $f_{BB}^S + f_{WW}^S$ . For large values the di-photon branching ratio of the heavy scalar can exceed 50%, while the remaining decay channel modes are suppressed, allowing  $\sin \alpha$  to increase without conflicting with data. If we allow for extreme values of  $f_{BB}^S/\Lambda + f_{WW}^S/\Lambda \sim 250 \text{ TeV}^{-1}$ , the upper bound on  $\sin \alpha$  goes up to 0.3. In the lower-center panel we again observe two distinct regions in  $f_{GG}^S$ . The vertical region with  $f_{GG}^S/\Lambda < 1.5 \text{ TeV}^{-1}$  is characterized by a large branching ratio for  $S \rightarrow \gamma\gamma$ , linked to large values of  $f_{BB}^S + f_{WW}^S$ . The horizontal region with  $f_{GG}^S/\Lambda = 1.5 \dots 10 \text{ TeV}^{-1}$  is characterized by a large production rate for the new scalar and a total decay width driven by  $f_{GG}^S$ . The upper limit on  $f_{GG}^S$  is set by di-jet searches, and the mixing in this regime has to be small to respect the limits from other decay channels. Finally, in the lower-right panel we show the correlation between the mixing angle and  $c_{SHH}$  from the limit on the decay  $S \rightarrow HH$ . Fixing  $c_{SHH} = 0$  and generating the  $SHH$  interaction through the mixing angle alone has no effect on any of the other correlations.

We proceed with an analysis allowing the new scalar to couple to the two fermions for which there are direct searches available. The analysis now includes

$$\{ f_{WW}^S, f_{BB}^S, f_{GG}^S, \sin \alpha, c_{SHH}, f_t^S, f_\tau^S \}. \quad (6)$$

A selection of results is shown in Fig. 2. The fermionic Wilson coefficients  $f_t^S$  and  $f_\tau^S$  are constrained

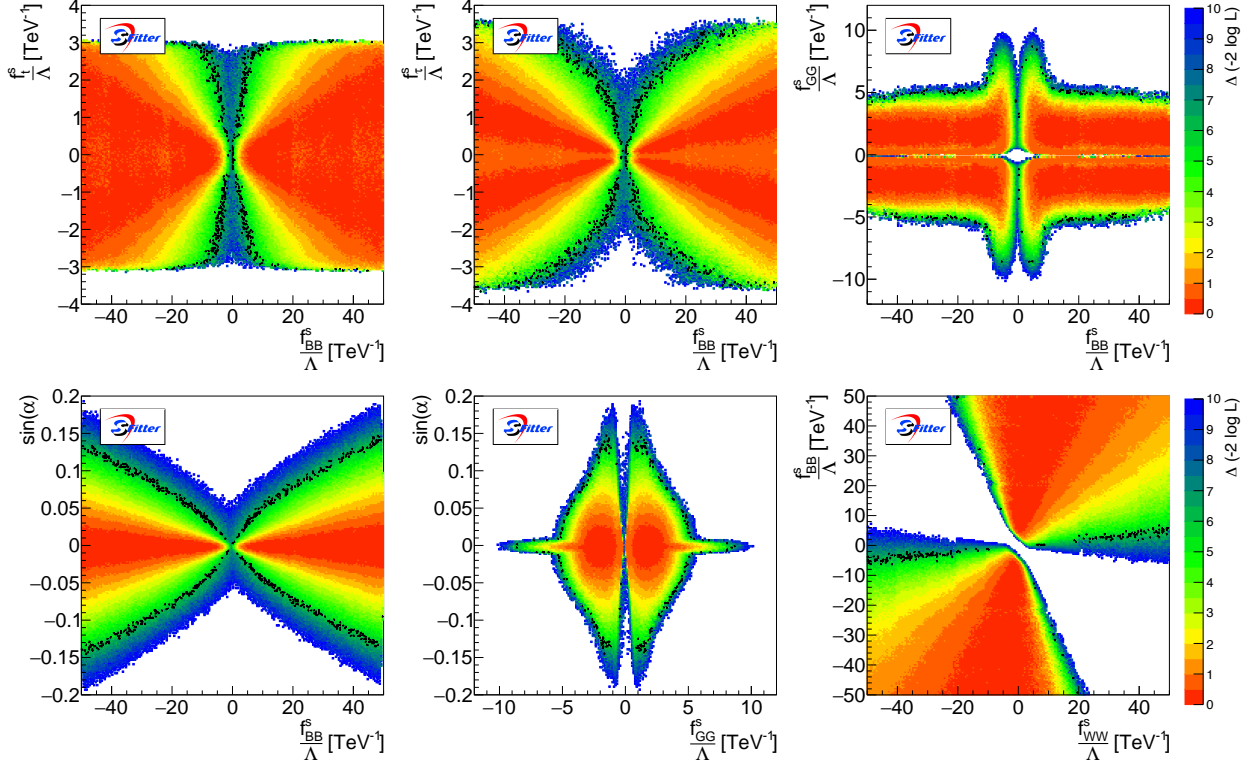


Figure 2. Two-dimensional profile log-likelihoods for the analysis of the heavy scalar sector alone. In contrast to Fig. 1 we now include fermion couplings in our set of Wilson coefficients  $f_{WW}^S$ ,  $f_{BB}^S$ ,  $f_{GG}^S$ ,  $\sin \alpha$ ,  $c_{SHH}$ ,  $f_t^S$  and  $f_\tau^S$ . The black points indicate  $\Delta(-2 \log L) = 5.99$ .

by  $t\bar{t}$  and  $\tau^+\tau^-$  resonance searches, as well as from an upper limit  $\Gamma_S < 25$  GeV which we assume throughout our analysis and which sets hard limits on  $f_t^S$  and  $f_\tau^S$ . The best fit point of this run is only mildly better than before,  $-2 \log L = 8.3$ . The limits on these two fermion couplings are stronger for smaller  $f_{BB}^S$ , as illustrated for  $f_t^S$  in the upper-left panel, and  $f_\tau^S$  in the upper-center one. The reason is that in those regions the partial decay width to photons becomes small, and the required di-photon branching ratio translates into small fermionic couplings. Conversely, larger fermionic Wilson coefficients now allow for best fit regions with large  $f_{GG}^S$  and  $f_{BB}^S$  at the same time, as shown in the upper-right panel of Fig. 2. This is the main difference with respect to the reduced analysis shown in Fig. 1. The rest of correlations remain qualitatively unchanged, as can be seen in the lower panels of Fig. 2. In particular the upper 95% CL limit on the mixing angle is still  $\sin \alpha < 0.15$ .

In passing we note that all the results shown so far have been derived assuming a CP-even new scalar. Nevertheless, for the analysis up to this point the results remain unchanged when instead we assume a heavy CP-odd scalar.

## B. Combined Higgs portal fit

Next, we discuss the results for the general scenario, where we constrain the 17 parameters

$$\{f_{WW}^S, f_{BB}^S, f_{GG}^S, \sin \alpha, f_t^S, f_b^S, f_\tau^S, f_{WW}, f_{BB}, f_{GG}, f_W, f_B, f_{\phi,2}, f_{WWW}, f_t, f_b, f_\tau\} \quad (7)$$

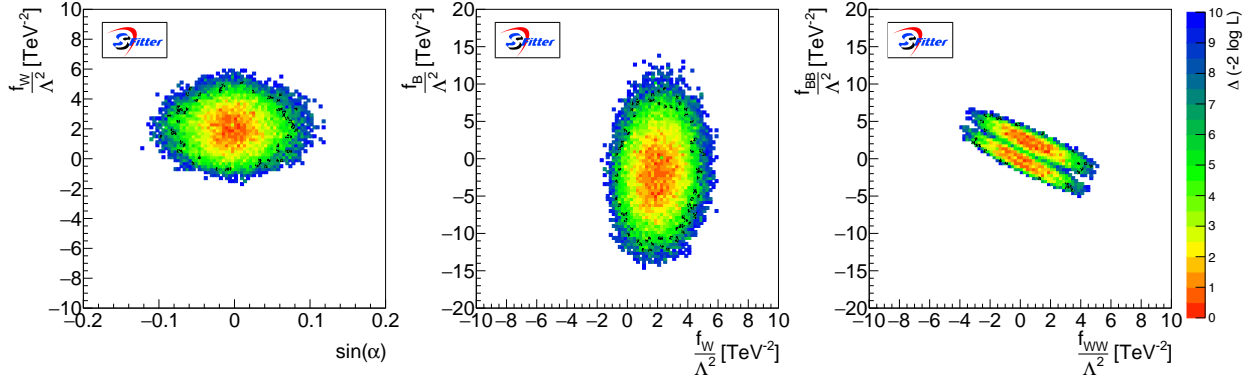


Figure 3. Two-dimensional profile log-likelihoods for the combined Higgs, TGV, and heavy scalar sectors. The black points indicate  $\Delta(-2 \log L) = 5.99$ .

from the combined measurements in the electroweak-Higgs, and the heavy scalar sector. We have fixed  $c_{SHH} = 0$  given its minor impact on the fit results.

In this case the best fit value has a likelihood of  $-2 \log L = 242.0$ , for an analysis containing 252 measurements, while the Standard Model point leads to  $-2 \log L = 273.9$ . In Fig. 3 we show a reduced selection of correlations between Wilson coefficients. When adding the heavy scalar to the combined Higgs and gauge boson analysis, the potentially largest change in the results appears for  $f_W$  and  $f_B$ . The twofold reason is illustrated in detail in the Appendix. First, focusing on the electroweak-Higgs phenomenology, while the contribution of  $f_W$  and  $f_B$  to the Higgs vertices is now weighted by the cosine of the mixing angle, their contribution to the triple gauge boson vertex is not. This generates a different pattern of Higgs-TGV correlations once we add the new scalar. Second, the mixing of the Higgs boson with the heavy scalar allows  $f_W$  and  $f_B$  to generate genuinely new Lorentz structure contributions to the  $SWW$ ,  $SZZ$  and  $SZ\gamma$  vertices, on top of the contributions from the rest of dimension-five and dimension-six operators.

The first effect turns out to be negligible, and given the small allowed size for the mixing angle, the electroweak-Higgs measurements are not precise enough to be sensitive to the scalar mixing contributions. Conversely, the second effect is more important. The mild preference for non-zero  $f_W$  values from the electroweak-Higgs measurements [5] causes the best fit regions to generate the new contribution to the decays  $S \rightarrow WW, ZZ, Z\gamma$ . These channels can be then better fit suppressing them further with a smaller mixing angle. The addition of the dimension-six operator causes then the upper bound on the scalar mixing to be mildly reduced with respect to the results in the previous section: now  $\sin \alpha < 0.10$  at 95% CL. This can be observed in the left panel of Fig. 3.

Apart from this effect, the small mixing angle causes a lack of sizable correlations between both the new scalar sector and the electroweak-Higgs sector. Consequently, the results and two-dimensional planes involving dimension-five operators are very similar to the ones shown in Fig. 2. The planes involving dimension-six operators remain unchanged with respect to the results shown in Ref. [5], something that we illustrate in the center and right panels of Fig. 3 for two of the dimension-six correlations.

### C. A common origin of operators

When we split a common scalar potential for two mixing states into a set of dimension-five and dimension-six operators, the question becomes how different the higher-dimensional effects in the



light and heavy scalar couplings can really be. In this section we assume that the set of heavy scalar couplings are directly tied to their Higgs-like counter parts,

$$\begin{aligned} \frac{f_{GG}}{\Lambda^2} &= -2 \frac{f_{GG}^S}{\Lambda} \left| \frac{f_{GG}^S}{\Lambda} \right| & \frac{f_f}{\Lambda^2} &= -\frac{v}{m_f} \frac{f_f^S}{\Lambda} \left| \frac{f_f^S}{\Lambda} \right| \\ \frac{f_{BB}}{\Lambda^2} &= -\frac{1}{4\pi^2} \frac{f_{BB}^S}{\Lambda} \left| \frac{f_{BB}^S}{\Lambda} \right| & \frac{f_{WW}}{\Lambda^2} &= -\frac{1}{4\pi^2} \frac{f_{WW}^S}{\Lambda} \left| \frac{f_{WW}^S}{\Lambda} \right|, \end{aligned} \quad (8)$$

for  $f = b, t, \tau$ . The relative signs and pre-factors ensure that the underlying new physics scales are consistent, as defined in Eq.(3). For the fermion case, this is motivated by the need to have a minimal flavor violating structure in both dimension-five and dimension-six operators to avoid large flavor changing neutral currents [59]. In a Bayesian language this approach would correspond to a Dirichlet prior, for example employed in the dark matter fit of Ref. [60], with an exponent parameter  $\alpha \gg 1$ .

After imposing the relations in Eq.(8), we proceed to perform the combined Higgs, triple gauge boson vertex and heavy scalar analysis spanning the 11 free parameters

$$\{ f_{WW}^S, f_{BB}^S, f_{GG}^S, f_t^S, f_b^S, f_\tau^S, \sin \alpha, f_W, f_B, f_{\phi,2}, f_{WW} \} . \quad (9)$$

We have again fixed  $c_{SHH} = 0$ , while  $f_{WW}$ ,  $f_{BB}$ ,  $f_{GG}$ ,  $f_t$ ,  $f_b$  and  $f_\tau$  are set from Eq.(8). Interestingly, the best fit point is  $-2 \log L = 242.6$ , i. e. within the analysis precision very close to the best fit point of the previous general scenario. This illustrates one of the most important conclusions: when dimension-five and dimension-six operators of a similar type are imposed to be related, there are still regions in the new physics parameter space which can accommodate the di-photon anomaly while respecting the constraints from the electroweak-Higgs measurements.

In Fig. 4 we again show a selection of two-dimensional correlations. In the upper-left panel we start with tight constraints on  $f_{BB}^S$  and also on  $f_{GG}^S$ . Now  $f_{BB}^S/\Lambda$  no longer presents an unconstrained direction, as the reduced allowed region for values around  $-10 \text{ TeV}^{-1}$  is limited from the constraint that  $f_{BB}$  and hence  $f_{BB}^S$  is constrained by the Higgs measurements. Because of the minus signs in Eq.(8) the region of allowed values for both  $f_{BB}^S$  and  $f_{WW}^S$  corresponds to the solution that flips the sign of the  $H\gamma\gamma$  vertex while respecting its measured size [5], as seen in the upper-center panel. In the case of  $f_{GG}^S$  and  $f_{GG}$  the several best fit regions are due to the measurement of a SM-like Higgs boson in gluon fusion production, the interference between  $f_{GG}$  and  $f_t$  [4], and the heavy scalar anomaly that excludes  $f_{GG}$  null values.

As seen in the upper-right panel, the stronger constraints on  $f_{BB}^S$  are directly translated into a stringent 95% CL bound on the mixing angle,  $\sin \alpha < 0.02$ . In the lower-left panel we show the impact of Eq.(8) on  $f_t^S$ . The fact that in this analysis  $f_{BB}^S$  is more constrained than in the general scenario implies that  $f_t^S$  is constrained to order one values, as expected from the  $f_{BB}^S$  vs  $f_t^S$  correlation in Fig. 2. The solution for  $f_t$  that flips the sign of the Higgs-Yukawa present in Ref. [4] is excluded through  $f_t^S$ . This reduces the number of allowed regions for  $f_{GG}$ , as compared to the electroweak-Higgs fit [4]. In the case of  $f_\tau^S$  and  $f_b^S$ , the allowed regions are limited by the  $H\tau^+\tau^-$  and  $Hb\bar{b}$  measurements. The  $v/m_f$  factors in Eq.(8) lead to reduced allowed ranges for  $f_\tau^S$  in comparison to the previous general scenario.

We illustrate in the lower-center panel the allowed region for two of the dimension-six operators not involved in the simplifications of Eq.(8),  $f_W$  and  $f_B$ . They remain unaltered with respect to the general analysis or the electroweak-Higgs results [5]. Conversely, in the lower-right panel of Fig. 4 we illustrate the two parameter regions  $f_{BB}$  vs.  $f_{WW}$ . There we see how the SM solution

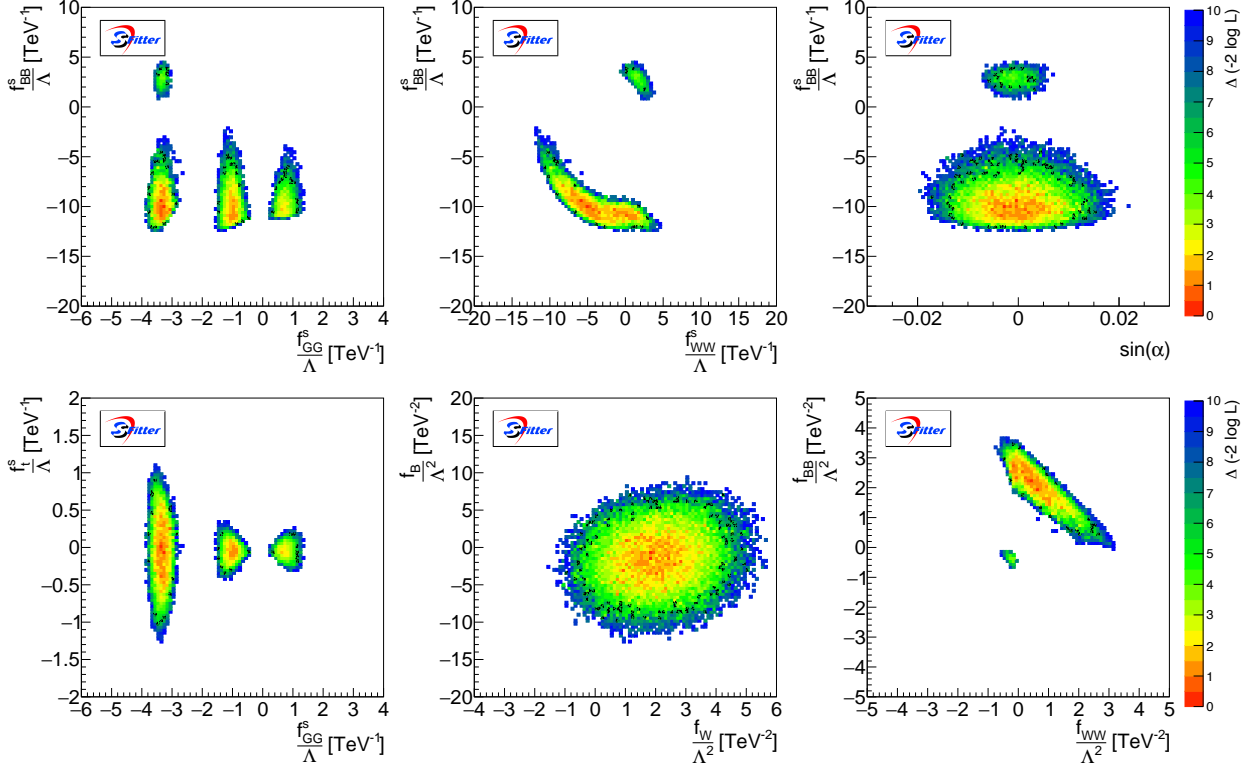


Figure 4. Two-dimensional profile log-likelihoods for the combined Higgs, TGV, and heavy scalar fit, but assuming a common origin of operators as defined in Eq.(8). The black points indicate  $\Delta(-2 \log L) = 5.99$ .

observed in the electroweak-Higgs analysis is now disfavored with respect to positive values for the Wilson coefficients.

In this section we have illustrated the results of a constrained scenario imposing hard relations between the heavy scalar and Higgs operators in Eq.(8). Realistically, we would expect such relations to not be as strict. We therefore checked that relaxing Eq.(8) and allowing for order-one variations does not qualitatively change our conclusions. Numerically, the bound on the mixing angle  $\sin \alpha$  becomes weaker once the relation  $f_{BB} \propto f_{BB}^S |f_{BB}^S|$  is relaxed.

### III. EPITAPH

We have developed the framework to perform a combined analysis of the electroweak-Higgs sector extended with a new scalar to test Higgs portal scenarios. The theoretical description we have studied is that of a linear effective Lagrangian extended with the addition of a singlet scalar.

The key question we face is the test of the portal structure hypothesis for an extended scalar sector. With that purpose we include a large set of Higgs event rates and kinematic distributions, combined with the recently implemented LHC triple gauge boson vertex distributions [5]. As a test of a Higgs portal scenario we study the possibility that a di-photon anomaly recently observed at the LHC [8–10] could be part of an extended Higgs sector. For that we include a selection of relevant experimental searches for heavy resonances as listed in Tab. I.

Analyzing first the new scalar sector only, we recover the result that a non-zero value for a reduced set of singlet scalar effective operators ( $f_{GG}^S, f_{BB}^S + f_{WW}^S$ ) fits the observed anomaly in

the di-photon channel, without conflicting with the lack of other positive observations, see Fig. 1. The mixing angle of the new singlet state with the Higgs boson can be sizable, the upper bound we find in the analysis is  $\sin(\alpha) < 0.15$  at the 95% CL. The addition of fermionic dimension-five operators increases the allowed parameter space regions for the bosonic operators. However it has no impact on the maximum allowed mixing angle value, see Fig. 2.

We then extend the analysis combining the new scalar sector with the electroweak-Higgs sector, using the Lagrangian description based on the dimension-six operators in Eq.(2). In this extended scenario the upper bound on the mixing angle is further reduced in order to suppress the new dimension-six contributions to the heavy scalar non-observed decays. The upper bound is now  $\sin(\alpha) < 0.1$  at 95% CL, with a size still compatible with Higgs portal hypothesis. Beyond this change, the maximum allowed mixing angle reduces the correlations between the Higgs-electroweak phenomenology and the hypothetical heavy scalar interactions. This leads to results that in the most general scenario are very similar to the ones of the individual Higgs-electroweak [5], and heavy scalar analysis, respectively.

Motivated by a scalar portal scenario we define and test a hypothesis for a unique origin of the dimension-five and dimension-six operators studied in the analysis. Imposing Eq.(8) we find new physics regions of parameters that fit the di-photon anomaly while being consistent with the lack of deviations measured on the electroweak-Higgs measurements. The upper bound on the mixing angle is reduced in this case to  $\sin(\alpha) < 0.02$ , due to the strong constraints on the operators modifying  $h \rightarrow \gamma\gamma$ .

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### HIGGS-SINGLET LAGRANGIAN

We describe here the main details of our effective Lagrangian analysis. We focus on the Higgs-scalar mixing and the combined phenomenology we derive. Following the Lagrangian in Eq.(3) of Sec. I A, both  $\mu_S$  and  $f_1^S/\Lambda$  generate a mixing between the two interaction eigenstates  $H'$  and  $S'$ . In this appendix we denote interaction eigenstates as primed fields, while mass eigenstates after the rotation

$$\mathcal{L}_m = -\frac{1}{2} (H' \ S') \begin{pmatrix} M_H^2 & v \left( \mu_S + \frac{f_1^S v^2}{\Lambda} \right) \left( 1 - \frac{f_{\phi,2} v^2}{2\Lambda^2} \right) \\ v \left( \mu_S + \frac{f_1^S v^2}{\Lambda} \right) \left( 1 - \frac{f_{\phi,2} v^2}{2\Lambda^2} \right) & M_S^2 + \frac{\lambda_{SH} v}{2} \end{pmatrix} \begin{pmatrix} H' \\ S' \end{pmatrix}, \quad (10)$$

as un-primed fields. The light mass term is  $M_H^2 = 2\lambda_H v^2(1 - v^2 f_{\phi,2}/\Lambda^2)$ , with the Higgs quartic  $\lambda$ . The contribution proportional from  $f_{\phi,2}/\Lambda^2$  is originated from the Higgs kinetic term and the appropriate field re-definition [3]. The physical masses are

$$M_{1,2}^2 = \frac{M_S^2 + \frac{\lambda_{SH} v}{2} + M_H^2}{2} \mp \frac{1}{2} \sqrt{\left( M_S^2 + \frac{\lambda_{SH} v}{2} - M_H^2 \right)^2 + 4v \left( \mu_S + \frac{f_1^S v^2}{\Lambda} \right)^2 \left( 1 - \frac{f_{\phi,2} v^2}{2\Lambda^2} \right)^2}, \quad (11)$$

and the mixing angle as a function of the physical masses reads

$$\sin 2\alpha = \frac{2v \left( \mu_S + \frac{f_1^S v^2}{\Lambda} \right) \left( 1 - \frac{f_{\phi,2} v^2}{2\Lambda^2} \right)}{M_2^2 - M_1^2} \xrightarrow{f_1^S=0} \mu_S = \sin 2\alpha \frac{M_2^2 - M_1^2}{2v} \left( 1 + \frac{v^2}{2} \frac{f_{\phi,2}}{\Lambda^2} \right). \quad (12)$$

The Higgs-scalar mixing affects many couplings of the mass eigenstates  $S$  and  $H$ . We first study the interactions of the light, Higgs-like, state. The admixture of the new scalar generates new interactions of the kind  $s_\alpha f_j^S / \Lambda$ , formally of dimension five, with an additional suppression by the mixing angle. Once we include the dimension-six operators of  $\mathcal{L}_{\text{dim-6}}^H$ , all mixing contributions can be absorbed in a re-definition of the effective Higgs Lagrangian, as long as we limit our analysis to tri-linear interactions. For example, the physical Higgs-gluon coupling becomes

$$g_{Hgg} = -\frac{\alpha_s}{8\pi} \left( c_\alpha \frac{f_{GG} v}{\Lambda^2} + 2s_\alpha \frac{f_{GG}^S}{\Lambda} \right) \equiv -\frac{\alpha_s}{8\pi} \frac{f'_{GG} v}{\Lambda^2} \quad (13)$$

where  $g_{Hgg}$  is defined through the term  $g_{Hgg} H G_{\mu\nu}^a G^{a\mu\nu}$  in the Lagrangian [4]. Using these kind of re-definitions the Higgs part of our analysis can be easily related to the results of Ref. [4, 5].

Because the Higgs-scalar mixing of Eq.(10) is defined in the broken phase and does not affect the Goldstone modes, this kind of re-definition does not apply to the triple gauge vertices constrained by di-boson production channels [4, 5]. The contribution of  $f_W$  and  $f_B$  in the Higgs sector is weighted by  $c_\alpha$ . For instance, the  $f_W$  contribution to the  $HWW$  interaction reads

$$\mathcal{L}^{HVV} \supset g_{HWW}^{(1)} (W_{\mu\nu}^+ W^{-\mu} \partial^\nu H + \text{h.c.}) \quad \text{with} \quad g_{HWW}^{(1)} = c_\alpha \frac{g^2 v}{2\Lambda^2} \frac{f_W}{2}. \quad (14)$$

In contrast, the contributions of  $f_W$  and  $f_B$  to the triple gauge boson vertices is not modified by such a mixing angle and remains the same as in the Higgs-gauge analysis [5]. This way, a sizable mixing with the heavy scalar changes the pattern of Higgs-TGV correlations.

On the heavy scalar side, the interaction with the incoming gluons is

$$g_{Sgg} = -\frac{\alpha_s}{8\pi} \left( s_\alpha \frac{f_{GG} v}{\Lambda^2} - 2c_\alpha \frac{f_{GG}^S}{\Lambda} \right) \equiv \frac{\alpha_s}{4\pi} \frac{f_{GG}^{S'}}{\Lambda}. \quad (15)$$

While the contributions of  $f_{WW} \leftrightarrow f_{WW}^S$ ,  $f_{BB} \leftrightarrow f_{BB}^S$  and the fermionic interactions  $f_f \leftrightarrow f_f^S$  follow this structure, the case of  $f_W$  and  $f_B$  is again special. Both Higgs-like operators generate new Lorentz structures in the heavy scalar sector. For example, the  $f_W$  contribution to the  $SWW$  vertex is

$$\mathcal{L}^{SVV} \supset g_{SWW}^{(1)} (W_{\mu\nu}^+ W^{-\mu} \partial^\nu S + \text{h.c.}) \quad \text{with} \quad g_{SWW}^{(1)} = s_\alpha \frac{g^2 v}{2\Lambda^2} \frac{f_W}{2}. \quad (16)$$

Finally, the  $SHH$  interaction is generated through the terms

$$\mathcal{L} \supset \lambda^{SHH} SHH + \frac{f_{\phi,2}}{\Lambda^2} s_\alpha c_\alpha^2 (S \partial^\mu H \partial_\mu H + 2H \partial^\mu S \partial_\mu H), \quad (17)$$

where the momentum-independent coupling is composed of several terms in Eq.(2) and Eq.(3)

$$\begin{aligned}
\lambda^{SHH} &= -3c_\alpha^2 s_\alpha v \lambda_H \left(1 - \frac{3f_{\phi,2}v^2}{2\Lambda^2}\right) \\
&\quad + \frac{1}{2} (2c_\alpha s_\alpha^2 - c_\alpha^3) \mu_S \left(1 - \frac{f_{\phi,2}v^2}{\Lambda^2}\right) + \frac{1}{2} (2c_\alpha^2 s_\alpha - s_\alpha^3) v \lambda_{SH} \left(1 - \frac{f_{\phi,2}v^2}{2\Lambda^2}\right) - 3c_\alpha s_\alpha^2 a_3 \\
&\quad + \frac{3}{4} (c_\alpha s_\alpha^2 - c_\alpha^3) \frac{f_1^S v^2}{\Lambda} \left(1 - \frac{f_{\phi,2}v^2}{\Lambda^2}\right) \\
&\equiv -3c_\alpha^2 s_\alpha v \lambda_H \left(1 - \frac{3f_{\phi,2}v^2}{2\Lambda^2}\right) + \frac{1}{2} (2c_\alpha s_\alpha^2 - c_\alpha^3) s_{2\alpha} \frac{M_2^2 - M_1^2}{v} \left(1 - \frac{f_{\phi,2}v^2}{2\Lambda^2}\right) + c_{SHH} . \quad (18)
\end{aligned}$$

There the Higgs self coupling  $\lambda_H$  can be expressed as

$$\lambda_H = \frac{s_\alpha^2 M_2^2 + c_\alpha^2 M_1^2}{2v^2} \left(1 + \frac{f_{\phi,2}v^2}{\Lambda^2}\right) , \quad (19)$$

while the term  $c_{SHH}$  accounts for the contributions from  $\lambda_{SH}$ ,  $a_3$ ,  $f_1^S$  and the dimension-6 operator  $(\phi^\dagger \phi)^3$ . In a simplified ansatz, we set  $c_{SHH} = 0$ .

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